Published for SISSA by Springer

Received: *March 5, 2013*
Accepted: *April 7, 2013*
Published: *April 29, 2013*

Higher spin black hole entropy in three dimensions

Alfredo Pérez,[a] **David Tempo**[a] **and Ricardo Troncoso**[a,b]

[a] *Centro de Estudios Científicos (CECs), Casilla 1469, Valdivia, Chile*
[b] *Universidad Andrés Bello, Av. República 440, Santiago, Chile*

E-mail: aperez@cecs.cl, tempo@cecs.cl, troncoso@cecs.cl

Abstract: A generic formula for the entropy of three-dimensional black holes endowed with a spin-3 field is found, which depends on the horizon area A and its spin-3 analogue $\varphi_+^{1/3}$, given by the reparametrization invariant integral of the induced spin-3 field at the spacelike section of the horizon. From this result it can be shown that the absolute value of φ_+ has to be bounded from above according to $|\varphi_+|^{1/3} \leq A/\sqrt{3}$. The entropy formula is constructed by requiring the first law of thermodynamics to be fulfilled in terms of the global charges obtained through the canonical formalism. For the static case, in the weak spin-3 field limit, our expression for the entropy reduces to the result found by Campoleoni, Fredenhagen, Pfenninger and Theisen, which has been recently obtained through a different approach.

Keywords: Black Holes, Classical Theories of Gravity, Gauge-gravity correspondence

ArXiv ePrint: 1301.0847

doi:10.1007/JHEP04(2013)143

Contents

1 Introduction

The successful formulation of a consistent classical field theory that describes interacting massless fields of spin greater than two [1–3] has recently attracted a great deal of attention, mainly coming from different research branches within the theoretical high energy physics community (for recent reviews see e.g., [4–9]). The three-dimensional case has been shown to be particularly useful in order to gather more intuition about this subject, in which spacetime and gauge symmetries become intrinsically mixed in an unfamiliar way. Indeed, in three spacetime dimensions it is possible to formulate a simpler theory that contains non-propagating but interacting fields of spins two and three only, without the need of introducing the whole tower of higher spin fields [10–13]. The Lagrangian is given by a Chern-Simons form constructed out from two copies of $SL(3,\mathbb{R})$, so that the theory describes gravity with negative cosmological constant, nonminimally coupled to an interacting spin-three field. The remarkable simplicity of this theory further allows to find exact black hole solutions endowed with a nontrivial spin-three field, as it has been recently reported in [14–16]. Note that, since the spin-three field is nonminimally coupled to gravity, one should not expect the standard Bekenstein-Hawking formula for the black hole entropy to apply for this class of higher spin black holes. In fact, in the simpler case of black holes with nonminimally coupled scalar hair, it is known that the entropy not only depends on the area, but also on the value of the matter field at the horizon [17–20]. It is then natural to think that a similar effect should occur in the case of higher spin black holes. One of the purposes of this paper is to show that this is indeed the case. For the theory under consideration, the higher spin black hole entropy is found to be given by

$$S = \frac{A}{4G} \cos\left[\frac{1}{3}\arcsin\left(3^{3/2}\frac{\varphi_+}{A^3}\right)\right], \tag{1.1}$$

where A and $\varphi_+^{1/3}$ stand for the reparametrization invariant integrals of the induced metric and spin-3 field at the spacelike section of the horizon, respectively; i.e., the horizon area

$$A = \int_{\partial\Sigma_+} \left(g_{\mu\nu}\frac{dx^\mu}{d\sigma}\frac{dx^\nu}{d\sigma}\right)^{1/2} d\sigma\,, \tag{1.2}$$

and its spin-three analogue

$$\varphi_+^{1/3} := \int_{\partial\Sigma_+} \left(\varphi_{\mu\nu\rho}\frac{dx^\mu}{d\sigma}\frac{dx^\nu}{d\sigma}\frac{dx^\rho}{d\sigma}\right)^{1/3} d\sigma\,. \tag{1.3}$$

Note that the presence of the spin-3 field does not spoil the Bekenstein bound, since the higher spin black hole entropy (1.1) is bounded as $S \leq A/(4G)$. Moreover, consistency of eq. (1.1) implies that the absolute value of φ_+ also has to be bounded from above according to

$$|\varphi_+|^{1/3} \leq \frac{A}{\sqrt{3}}\,. \tag{1.4}$$

The expression for the higher spin black hole entropy in (1.1) can be readily obtained by virtue of the first law of thermodynamics. In order to perform this task, one short and simple path turns out to be working in the canonical ensemble, since in this case, apart from the Hawking temperature, one just need to know the variation of the total energy of the system. In this way one circumvents the explicit computation of higher spin charges and their corresponding chemical potentials that contribute to the work terms in the grand canonical ensemble. As it was shown in [21], the canonical approach of [22] applied to the Chern-Simons theory for $SL(3,\mathbb{R})\times SL(3,\mathbb{R})$ (see also e.g. [23–25]), allows to obtain the variation of the total energy as the surface integral that corresponds to the variation of the total Hamiltonian, which is given by

$$\delta E = \delta Q(\partial_t) = \frac{k}{2\pi}\int_{\partial\Sigma}\left(\langle A_t^+\delta A_\theta^+\rangle - \langle A_t^-\delta A_\theta^-\rangle\right)d\theta\,, \tag{1.5}$$

where $\partial\Sigma$ stands for the boundary of the spacelike section. The level is determined by the Newton constant and the AdS radius according to $k=\frac{l}{4G}$, while $A_\mu^\pm$ correspond to the gauge fields associated to both copies of $SL(3,\mathbb{R})$. Here the bracket $\langle\cdots\rangle$ is given by a quarter of the trace in the fundamental representation (see e.g., [13]).

Since (1.5) corresponds to the variation of the energy in the canonical ensemble, the higher spin black hole entropy (1.1) can be found from the first law $\delta E = T\delta S$, where T stands for the Hawking temperature. In what follows an explicit check of this approach is carried out for the higher spin black hole solutions found in [14, 15] and [16].[1]

2 Explicit examples

(i) The exact solution that has been found in [14, 15] is described by

$$A^\pm = g_\pm^{-1}a^\pm g_\pm + g_\pm^{-1}dg_\pm\,, \tag{2.1}$$

where $g_\pm = g_\pm(\rho)$ correspond to suitable group elements of each copy of $SL(3,\mathbb{R})$ that depend only on the radial coordinate, and

$$\begin{aligned} a^\pm = &\pm\left(L_{\pm1}^\pm - \frac{2\pi}{k}\mathcal{L}L_{\mp1}^\pm \mp \frac{\pi}{2k}\mathcal{W}W_{\mp2}^\pm\right)dx^\pm \\ &+\mu\left(W_{\pm2}^\pm - \frac{4\pi}{k}\mathcal{L}W_0^\pm + \frac{4\pi^2}{k^2}\mathcal{L}^2W_{\mp2}^\pm \pm \frac{4\pi}{k}\mathcal{W}L_{\mp1}^\pm\right)dx^\mp, \end{aligned} \tag{2.2}$$

[1] Different aspects of the thermodynamics of these solutions have been discussed in [26–28].

where $x^{\pm} = \frac{1}{l}t \pm \theta$. This solution describes a static higher spin black hole, whose metric asymptotically approaches to that of AdS_3 of radius $\tilde{l} = l/2$, so that the asymptotic behaviour is relaxed as compared with the one of [13, 29]. As shown in [21], the explicit form of $g_{\pm} = g_{\pm}(\rho)$, is not needed to compute the variation of energy, since (1.5) reduces to

$$\delta E = \delta Q(\partial_t) = \frac{k}{2\pi} \int_{\partial\Sigma} \left(\langle a_t^+ \delta a_\theta^+ \rangle - \langle a_t^- \delta a_\theta^- \rangle\right) d\theta \,. \tag{2.3}$$

Therefore, evaluating (2.3) for the gauge fields $a^{\pm}$ in (2.2), the variation of the energy was found to be given by

$$\delta E = \frac{8\pi}{l} \left[\delta\mathcal{L} - \frac{32\pi}{3k} \delta(\mathcal{L}^2 \mu^2) + \mu \delta\mathcal{W} + 3\mathcal{W}\delta\mu \right] . \tag{2.4}$$

As explained in [15], it is useful to introduce the following change of variables

$$\frac{C-1}{C^{3/2}} = \sqrt{\frac{k}{32\pi\mathcal{L}^3}} \mathcal{W} \, ; \qquad \gamma = \sqrt{\frac{2\pi\mathcal{L}}{k}} \mu \, , \tag{2.5}$$

since the conditions that are obtained from requiring the holonomy along the thermal cycle for the Euclidean solution to be trivial become fairly simplified. One of these conditions allows to find the Hawking temperature $T = \beta^{-1}$, with

$$\beta = \frac{l}{4} \sqrt{\frac{\pi k}{2\mathcal{L}}} \frac{(2C-3)}{(C-3)} \left(1 - \frac{3}{4C}\right)^{-\frac{1}{2}} , \tag{2.6}$$

where in eq. (2.6) we have made use of the remaining condition.[2] The variation of the energy (2.4) then reads

$$\delta E = \frac{8\pi}{l} \frac{(C-3)}{(2C-3)^2} \left[(4C-3)\delta\mathcal{L} - \frac{9}{C} \frac{(2C-1)}{(2C-3)} \mathcal{L} \delta C \right] . \tag{2.7}$$

Note that the variation of the canonical energy is not an exact differential. Nevertheless, as it has to be in the canonical ensemble, the inverse Hawking temperature β acts as an integrating factor such that the product $\beta\delta E$ becomes an exact differential that defines the variation of the entropy, i.e.,

$$\delta S = \beta \delta E = \delta \left[4\pi \sqrt{2\pi k \mathcal{L}} \left(1 - \frac{3}{2C}\right)^{-1} \sqrt{1 - \frac{3}{4C}} \right] , \tag{2.8}$$

and therefore, the entropy is given by

$$S = 4\pi \sqrt{2\pi k \mathcal{L}} \left(1 - \frac{3}{2C}\right)^{-1} \sqrt{1 - \frac{3}{4C}} \, . \tag{2.9}$$

[2] The inverse Hawking temperature in (2.6) differs from the one in [14, 15] by a factor 1/2, since here we are using the time scale of the asymptotic region that corresponds to AdS_3 spacetime of radius $\tilde{l} = l/2$.

In terms of the horizon area and the value of the purely angular component of the spin-3 field at the horizon, which for this case read [15][3]

$$g_{\theta\theta}|_{\rho_+} = \left(\frac{A}{2\pi}\right)^2 = \frac{8\pi l^2}{k}\frac{4C^3+9C^2-36C+27}{C(2C-3)^2}\mathcal{L}\,, \tag{2.10}$$

$$\varphi_{\theta\theta\theta}|_{\rho_+} = \frac{\varphi_+}{(2\pi)^3} = -\left(\frac{8\pi l^2\mathcal{L}}{kC}\right)^{3/2}\frac{(C-1)(4C^2-3C+9)}{(2C-3)^2}\,, \tag{2.11}$$

one verifies that the entropy formula in eq. (1.1) reduces to (2.9).

It is worth pointing out that the perturbative expansion of the higher spin black hole entropy in the spin-three field, in terms of the original variables $\mathcal{L}$ and $\mathcal{W}$ reads

$$S = 4\pi\sqrt{2\pi k\mathcal{L}}\left[1+\frac{9k}{256\pi}\left(\frac{\mathcal{W}^2}{\mathcal{L}^3}\right)+\mathcal{O}\left(\frac{\mathcal{W}^2}{\mathcal{L}^3}\right)^2\right]\,, \tag{2.12}$$

in agreement with the result found in [30] through a different approach. The full nonperturbative expression for the entropy S in eq. (2.9) differs from the one found in [15], here denoted as $\tilde{S}$, by a factor that depends on the integration constant C that characterizes the spin-three field, i.e., $S = \tilde{S}\left(1-\frac{3}{2C}\right)^{-1}$. Further comments about the comparison with the results in [30] and [15] are discussed below.

(ii) In the case of the static higher spin black hole solution found in [16], whose metric approaches to that of AdS$_3$ of radius $\tilde{l} = l/2$, the gauge fields can also be expressed as in (2.1), with $g_\pm = e^{\pm\rho L_0^\pm}$ and

$$a^\pm = \pm\left(\ell_P L_{\pm1}^\pm - \mathcal{L}L_{\mp1}^\pm \pm \Phi W_0^\pm\right)dx^\pm + \left(\ell_D W_{\pm2}^\pm + \mathcal{W}W_{\mp2}^\pm - QW_0^\pm\right)dx^\mp, \tag{2.13}$$

with[4]

$$Q\ell_P - 2\mathcal{L}\ell_D = 0\,; \qquad Q\mathcal{L} - 2\mathcal{W}\ell_P = 0\,. \tag{2.14}$$

Requiring the holonomy around the thermal cycle of the Euclidean solution to be trivial, fixes the Hawking temperature as

$$T = \frac{2}{\pi l}\sqrt{\mathcal{L}\ell_p + 4Q^2}\,, \tag{2.15}$$

and also gives the condition $\Phi = 4Q$.

In [21], by making use of the surface integral in eq. (1.5), the total energy of this higher spin black hole was found to be given by

$$E = \frac{1}{G}[\mathcal{L}\ell_P + 4Q^2] = \frac{\pi^2 l^2}{4G}T^2, \tag{2.16}$$

so that by virtue of the first law $\delta E = T\delta S$, the entropy was shown to agree with Cardy formula, i.e.,

$$S = \pi l\sqrt{\frac{E}{G}} = 4\pi\sqrt{\frac{c}{6}L_0}\,, \tag{2.17}$$

[3] Here, the normalization of the spin-3 field agrees with [30], which differs by a factor 1/6 as compared with the one in [15].

[4] The orientability we have chosen differs from the one in [16], i.e., $x^+ \leftrightarrow x^-$.

JHEP04(2013)143

where, since the asymptotic form of the metric approaches to AdS_3 of radius $\tilde{l}$, the zero modes of the Virasoro generators are given by

$$L_0 := L_0^{\pm} = \frac{\tilde{l}}{2} E = \frac{l}{4} E\,, \tag{2.18}$$

and c stands for the central charge that was shown to agree with the one of Brown and Henneaux in [13, 29, 31], i.e., $c = 3l/(2G)$.

In this case, the value of the purely angular component of the spin-3 field at the horizon, and the horizon area are given by

$$\varphi_{\theta\theta\theta}|_{\rho_+} = \frac{\varphi_+}{(2\pi)^3} = -\frac{32}{3} l^3 Q \left(\mathcal{L}\ell_p + \frac{20}{9} Q^2 \right), \tag{2.19}$$

$$g_{\theta\theta}|_{\rho_+} = \left(\frac{A}{2\pi} \right)^2 = 4l^2 \left(\mathcal{L}\ell_p + \frac{28}{3} Q^2 \right), \tag{2.20}$$

respectively, from which it can seen that the entropy in eq. (2.17) for this higher spin black hole is also successfully reproduced by virtue of the generic entropy formula in (1.1).

3 Discussion and final remarks

It has been shown that the entropy of black holes endowed with a spin-3 field in three dimensions can be expressed in terms of the reparametrization invariant integrals of the pullback of the metric and the spin-3 field at the spacelike section of the horizon, namely, the horizon area A, and its spin-3 analogue $\varphi_+^{1/3}$. The entropy formula (1.1) turns out to be, by construction, invariant under proper gauge transformations. Indeed, since the variation of the total energy (1.5) is invariant by definition, as well as it is the temperature, because it can be found through requiring the holonomy along the thermal cycle to be trivial, by virtue of the first law, $\delta S = \beta \delta E$, the entropy does.

Consistency of (1.1) also implies that the absolute value of $\varphi_+^{1/3}$ becomes bounded from above by the horizon area, as in eq. (1.4). It is then worth pointing out that for the allowed range of parameters for which the higher spin black holes of refs. [14, 15] and [16] are defined, eq. (1.4) turns out to be at most saturated, which reassuringly means that the bound cannot be improved.

Note that as expressed by formula (1.1), the higher spin black hole entropy turns out to be described by a multivalued function, so that in order to suitably take into account the different branches of the codomain, it is useful to parametrize the quotient φ_+/A^3 by an "angular" variable, defined as

$$\sin\left(\frac{\phi}{2}\right) := 3^{3/2} \frac{\varphi_+}{A^3}\,, \tag{3.1}$$

so that the entropy reads

$$S = \frac{A}{4G} \cos\left(\frac{\phi}{6}\right). \tag{3.2}$$

JHEP04(2013)143

Positivity of the entropy then implies that the angle fulfills $|\phi| \leq 3\pi$. In terms of this variable, the bound in (1.4) becomes saturated for $|\phi| = \pi$ or $|\phi| = 3\pi$, so that and in these cases the entropy can be either $S = \frac{\sqrt{3}}{2}\frac{A}{4G}$, or $S = 0$, respectively. It is also worth highlighting that φ_+ may vanish even in the presence of a nontrivial spin-3 field along spacetime, which corresponds to the cases $\phi = 0$, and $|\phi| = 2\pi$. In the former case (no winding) the entropy reduces to the one of a BTZ black hole constructed out from identifications of the AdS_3 vacuum of radius l, i.e., $S = A/(4G)$, while for the latter (winding one), the entropy is given by $S = A/(8G)$, which also corresponds to the one of a BTZ black hole, but constructed out from the other maximally symmetric vacuum, described by AdS_3 spacetime of radius $\tilde{l} = l/2$.

In the case of static circularly symmetric higher spin black holes, since the metric can be written in diagonal form, the horizon is located at a fixed value of the radial coordinate $\rho = \rho_+$, so that the cube of the spin-3 analogue of the area becomes determined by the purely angular component of the spin-3 field at the horizon according to $\varphi_+ = (2\pi)^3 \varphi_{\theta\theta\theta}|_{\rho_+}$. Hence, the entropy (1.1) reduces to

$$S = \frac{A}{4G}\cos\left[\frac{1}{3}\arcsin\left(3^{3/2}\frac{\varphi_{\theta\theta\theta}}{g_{\theta\theta}^{3/2}}\right)\right]\Bigg|_{\rho_+}, \tag{3.3}$$

where $A = 2\pi\sqrt{g_{\theta\theta}}|_{\rho_+}$. It is then reassuring to verify that in the weak spin-3 field limit, formula (3.3) expands as

$$S = \frac{A}{4G}\left(1 - \frac{3}{2}\left(g^{\theta\theta}\right)^3 \varphi_{\theta\theta\theta}^2 + \mathcal{O}(\varphi^4)\right)\Bigg|_{\rho_+}, \tag{3.4}$$

in full agreement with the result found by Campoleoni, Fredenhagen, Pfenninger and Theisen [30], which has been recently obtained through a completely different approach. Indeed, in ref. [30], the action was explicitly expressed in terms of the metric and the perturbative expansion of the spin-3 field up to quadratic order, so that the correction to the area law described by (3.4) was found by means of Wald's formula [32].

The discrepancy of our results for the higher spin black hole entropy, as compared with the ones in [15],[5] stems from a mismatch in the global charges that becomes inherited by the entropy once computed through the first law. In fact, as explained in [21], since the higher spin black hole solution found in [14, 15] possesses a relaxed asymptotic behavior with respect to the one described in refs. [13, 29, 31], the global charges, and in particular, the variation of the total energy (1.5), do not depend linearly on the deviation of the fields with respect to the reference background, so that the additional nonlinear contributions appearing in (2.4) cannot be neglected even in the weak spin-3 field limit.

As an ending remark, it would be interesting to explore how the higher spin black hole entropy formula (1.1) and the bound (1.4) extend to the case of theories that include fields with spin greater than three, as it is the case for the solutions discussed in [33, 35–37], as well as how could they be recovered, along the lines of [33, 34, 38–43], from a suitable notion of energy for the dual theory at the boundary.

[5] For an interpretation of the proposal in [15], in terms of an underlying dual CFT in two dimensions, see also [33, 34].

Acknowledgments

We thank G. Barnich, A. Campoleoni, G. Compere, C. Martínez, M. Henneaux and C. Troessaert, for useful discussions and comments. R.T. thanks the Physique théorique et mathématique group of the Université Libre de Bruxelles and the International Solvay Institutes for the kind hospitality, and the Instituut voor Theoretische Fysica, Katholieke Universiteit Leuven, for the opportunity of presenting this work. This research has been partially supported by Fondecyt grants N° 1095098, 1121031, 3110122, 3110141, and by the Conicyt grant ACT-91: "Southern Theoretical Physics Laboratory" (STPLab). The Centro de Estudios Científicos (CECs) is funded by the Chilean Government through the Centers of Excellence Base Financing Program of Conicyt.

References

[1] E.S. Fradkin and M.A. Vasiliev, *On the gravitational interaction of massless higher spin fields*, *Phys. Lett.* **B 189** (1987) 89 [INSPIRE].

[2] M.A. Vasiliev, *Consistent equation for interacting gauge fields of all spins in (3+1)-dimensions*, *Phys. Lett.* **B 243** (1990) 378 [INSPIRE].

[3] M.A. Vasiliev, *Nonlinear equations for symmetric massless higher spin fields in (A)dS$_d$*, *Phys. Lett.* **B 567** (2003) 139 [hep-th/0304049] [INSPIRE].

[4] M.A. Vasiliev, *Higher spin gauge theories in four-dimensions, three-dimensions and two-dimensions*, *Int. J. Mod. Phys.* **D 5** (1996) 763 [hep-th/9611024] [INSPIRE].

[5] X. Bekaert, S. Cnockaert, C. Iazeolla and M.A. Vasiliev, *Nonlinear higher spin theories in various dimensions*, hep-th/0503128 [INSPIRE].

[6] X. Bekaert, N. Boulanger and P. Sundell, *How higher-spin gravity surpasses the spin two barrier: no-go theorems versus yes-go examples*, *Rev. Mod. Phys.* **84** (2012) 987 [arXiv:1007.0435] [INSPIRE].

[7] A. Sagnotti, *Notes on strings and higher spins*, arXiv:1112.4285 [INSPIRE].

[8] M.R. Gaberdiel and R. Gopakumar, *Minimal model holography*, arXiv:1207.6697 [INSPIRE].

[9] M. Ammon, M. Gutperle, P. Kraus and E. Perlmutter, *Black holes in three dimensional higher spin gravity: a review*, arXiv:1208.5182 [INSPIRE].

[10] C. Aragone and S. Deser, *Hypersymmetry in $D = 3$ of coupled gravity-massless spin-5/2 system*, *Class. Quant. Grav.* **1** (1984) L9 [INSPIRE].

[11] M.P. Blencowe, *A consistent interacting massless higher-spin field theory in $D = (2 + 1)$*, *Class. Quant. Grav.* **6** (1989) 443 [INSPIRE].

[12] E. Bergshoeff, M.P. Blencowe and K.S. Stelle, *Area preserving diffeomorphisms and higher spin algebra*, *Commun. Math. Phys.* **128** (1990) 213 [INSPIRE].

[13] A. Campoleoni, S. Fredenhagen, S. Pfenninger and S. Theisen, *Asymptotic symmetries of three-dimensional gravity coupled to higher-spin fields*, *JHEP* **11** (2010) 007 [arXiv:1008.4744] [INSPIRE].

[14] M. Gutperle and P. Kraus, *Higher spin black holes*, *JHEP* **05** (2011) 022 [arXiv:1103.4304] [INSPIRE].

[15] M. Ammon, M. Gutperle, P. Kraus and E. Perlmutter, *Spacetime geometry in higher spin gravity*, *JHEP* **10** (2011) 053 [arXiv:1106.4788] [INSPIRE].

[16] A. Castro, E. Hijano, A. Lepage-Jutier and A. Maloney, *Black holes and singularity resolution in higher spin gravity*, *JHEP* **01** (2012) 031 [arXiv:1110.4117] [INSPIRE].

[17] M. Natsuume, T. Okamura and M. Sato, *Three-dimensional gravity with conformal scalar and asymptotic Virasoro algebra*, *Phys. Rev.* **D 61** (2000) 104005 [hep-th/9910105] [INSPIRE].

[18] M. Henneaux, C. Martínez, R. Troncoso and J. Zanelli, *Black holes and asymptotics of 2+1 gravity coupled to a scalar field*, *Phys. Rev.* **D 65** (2002) 104007 [hep-th/0201170] [INSPIRE].

[19] A. Ashtekar, A. Corichi and D. Sudarsky, *Nonminimally coupled scalar fields and isolated horizons*, *Class. Quant. Grav.* **20** (2003) 3413 [gr-qc/0305044] [INSPIRE].

[20] C. Martínez, J.P. Staforelli and R. Troncoso, *Topological black holes dressed with a conformally coupled scalar field and electric charge*, *Phys. Rev.* **D 74** (2006) 044028 [hep-th/0512022] [INSPIRE].

[21] A. Pérez, D. Tempo and R. Troncoso, *Higher spin gravity in 3D: black holes, global charges and thermodynamics*, arXiv:1207.2844 [INSPIRE].

[22] T. Regge and C. Teitelboim, *Role of surface integrals in the Hamiltonian formulation of general relativity*, *Annals Phys.* **88** (1974) 286 [INSPIRE].

[23] A.P. Balachandran, G. Bimonte, K.S. Gupta and A. Stern, *Conformal edge currents in Chern-Simons theories*, *Int. J. Mod. Phys.* **A 7** (1992) 4655 [hep-th/9110072] [INSPIRE].

[24] M. Bañados, *Global charges in Chern-Simons field theory and the (2+1) black hole*, *Phys. Rev.* **D 52** (1996) 5816 [hep-th/9405171] [INSPIRE].

[25] S. Carlip, *Conformal field theory, (2+1)-dimensional gravity and the BTZ black hole*, *Class. Quant. Grav.* **22** (2005) R85 [gr-qc/0503022] [INSPIRE].

[26] M. Bañados, R. Canto and S. Theisen, *The action for higher spin black holes in three dimensions*, *JHEP* **07** (2012) 147 [arXiv:1204.5105] [INSPIRE].

[27] J.R. David, M. Ferlaino and S.P. Kumar, *Thermodynamics of higher spin black holes in 3D*, *JHEP* **11** (2012) 135 [arXiv:1210.0284] [INSPIRE].

[28] B. Chen, J. Long and Y.-N. Wang, *Phase structure of higher spin black hole*, *JHEP* **03** (2013) 017 [arXiv:1212.6593] [INSPIRE].

[29] M. Henneaux and S.-J. Rey, *Nonlinear W_∞ as asymptotic symmetry of three-dimensional higher spin anti-de Sitter gravity*, *JHEP* **12** (2010) 007 [arXiv:1008.4579] [INSPIRE].

[30] A. Campoleoni, S. Fredenhagen, S. Pfenninger and S. Theisen, *Towards metric-like higher-spin gauge theories in three dimensions*, arXiv:1208.1851 [INSPIRE].

[31] A. Campoleoni, S. Fredenhagen and S. Pfenninger, *Asymptotic W-symmetries in three-dimensional higher-spin gauge theories*, *JHEP* **09** (2011) 113 [arXiv:1107.0290] [INSPIRE].

[32] R.M. Wald, *Black hole entropy is the Noether charge*, *Phys. Rev.* **D 48** (1993) 3427 [gr-qc/9307038] [INSPIRE].

[33] P. Kraus and E. Perlmutter, *Partition functions of higher spin black holes and their CFT duals*, *JHEP* **11** (2011) 061 [arXiv:1108.2567] [INSPIRE].

JHEP04(2013)143

[34] M.R. Gaberdiel, T. Hartman and K. Jin, *Higher spin black holes from CFT*, *JHEP* **04** (2012) 103 [arXiv:1203.0015] [INSPIRE].

[35] H.-S. Tan, *Aspects of three-dimensional spin-4 gravity*, *JHEP* **02** (2012) 035 [arXiv:1111.2834] [INSPIRE].

[36] M. Gary, D. Grumiller and R. Rashkov, *Towards non-AdS holography in 3-dimensional higher spin gravity*, *JHEP* **03** (2012) 022 [arXiv:1201.0013] [INSPIRE].

[37] B. Chen, J. Long and Y.-n. Wang, *Black holes in truncated higher spin AdS_3 gravity*, *JHEP* **12** (2012) 052 [arXiv:1209.6185] [INSPIRE].

[38] M.R. Gaberdiel and R. Gopakumar, *An AdS_3 dual for minimal model CFTs*, *Phys. Rev.* **D 83** (2011) 066007 [arXiv:1011.2986] [INSPIRE].

[39] M.R. Gaberdiel, R. Gopakumar, T. Hartman and S. Raju, *Partition functions of holographic minimal models*, *JHEP* **08** (2011) 077 [arXiv:1106.1897] [INSPIRE].

[40] M. Ammon, P. Kraus and E. Perlmutter, *Scalar fields and three-point functions in $D = 3$ higher spin gravity*, *JHEP* **07** (2012) 113 [arXiv:1111.3926] [INSPIRE].

[41] M.R. Gaberdiel and P. Suchanek, *Limits of minimal models and continuous orbifolds*, *JHEP* **03** (2012) 104 [arXiv:1112.1708] [INSPIRE].

[42] P. Kraus and E. Perlmutter, *Probing higher spin black holes*, *JHEP* **02** (2013) 096 [arXiv:1209.4937] [INSPIRE].

[43] S. Banerjee et al., *Smoothed transitions in higher spin AdS gravity*, arXiv:1209.5396 [INSPIRE].